

Selected problems in astrophysics of compact objects

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Abstract. I review three problems in astrophysics of compact stars: (i) the phase diagram of warm pair-correlated nuclear matter at sub-saturation densities at finite isospin asymmetry; (ii) the Standard Model neutrino emission from superfluid phases in neutron stars within the Landau theory of Fermi (superfluid) liquids; (iii) the beyond Standard Model physics of axionic cooling of compact stars by the Cooper pair-breaking processes.

1. Introduction

The present lecture summarizes some of the recent work on three problems in the astrophysics of compact stars listed in the abstract. It attempts to provide sufficient perspective and detail within the limited span of pages as to be a useful introduction to these problems. The selection of the topics is largely motivated by the author's interests and covers only a small part of many facets of compact star physics.

Neutron stars are born in spectacular explosions of type-II supernovas [1]. The equation of state of matter at subnuclear densities and more generally its phase structure and composition are important ingredients of supernova physics, one key reason being their importance for the formation of neutrino signal at the neutrino-sphere and for transport of the energy that ultimately may generate a successful explosion. The temperature range relevant for subnuclear densities is several MeVs. In Sec. 2 we will discuss some recent progress in the understanding of the phase diagram of nuclear matter in this regime.

If a neutron star is formed in a supernova explosion, its further thermal evolution is determined by the neutrino emission from the interior of the star: neutrinos once produced in the star escape the stellar environment without further interactions [2, 3, 4]. The black-body radiation from the surface of the star in thermal equilibrium, while providing us with the information on the composition of matter and effective temperature of the emitting regions, is an unimportant cooling agent until the star is hundreds of thousands years old. The physics of neutrino emission from the interiors of neutron stars is discussed in Sec. 3 in the case of (non-exotic) baryonic matter.

Non-Standard Model particles emitted during the neutrino stage of cooling of a neutron star can drain sufficient energy as to “spoil” the Standard Model based cooling scenario. This may be used to place limits on the properties of exotic particles, such as the axions. Astrophysical bounds on the properties of axions, i.e., the mass and the coupling to the Standard Model particles, have been derived from stellar (red giant, white dwarf, supernova, and neutron star) physics, which are complementary to those obtained from cosmology and terrestrial laboratory searches [5]. Section 4 discusses the bounds on the axion properties derived from the cooling of

neutron stars through inclusion of a new process of axion emission via the Cooper pair-breaking in superfluid baryonic matter [6].

2. Warm isospin asymmetrical nuclear matter

Nuclear matter at sub-saturation densities, i.e., at $\rho \leq 0.5 \rho_0$, where $\rho_0 = 2.8 \times 10^{14} \text{ g cm}^{-3}$ is the nuclear saturation density, is a many-body system with well defined pair interactions, whose dominant attractive part is responsible for the formation of nuclear clusters and the Bardeen-Cooper-Schrieffer (BCS) type pair-condensate(s) of nucleons. In the supernova context the matter is at finite, but small compared to cold neutron stars, isospin asymmetry. Furthermore, low densities and relatively high (again compared to neutron stars) temperatures allow for the existence of a substantial amount of light clusters, notably deuterons, tritons, He^3 nuclei and alpha particles [7, 8, 9, 10, 11, 12, 13].

Because the unbound nucleons and deuterons are the dominant component of the matter at all relevant densities it is useful to consider only two-body correlations first. In the extreme low density limit such matter can be considered as a mixture of quasi-free nucleons and deuterons, where the only effect of the interaction is to renormalize the mass of the constituents. The deuterons being bosons may also condense and form a Bose-Einstein condensate (BEC) even without interactions if, of course, the temperature is sufficiently low [14, 15, 16, 17, 18, 19]. Imagine now increasing density while keeping the temperature of the system constant, such that the degeneracy of the system is effectively increased. As the density $\rho \rightarrow \rho_0$ the abundances of deuterons are reduced because of the Pauli-blocking of the phase space available to nucleons [16, 20]. However, the emergence of the Fermi-surface of nucleons (at finite isospin - two Fermi surfaces of neutrons and of protons) leads to BCS type correlations and formation of macroscopic coherent state - the condensate of nucleons. The interaction channel responsible for this pairing phenomenon is the 3S_1 - 3D_1 partial wave, i.e., the same interaction channel which binds the deuteron. Thus, it has been conjectured that the nuclear matter may undergo a BCS-BEC phase transition, first in the context of intermediate energy heavy-ion collisions [14, 15] and more recently in the context of supernovas [8, 21].

The theoretical framework for the description of the BCS-BEC transition was developed by Nozières and Schmitt-Rink in condensed matter physics [22]. Isospin asymmetry, induced by weak interactions in stellar environments and expected in exotic nuclei, disrupts isoscalar neutron-proton (np) pairing, since the mismatch in the Fermi surfaces of protons and neutrons suppresses the pairing correlations [23]. The standard Nozières-Schmitt-Rink theory of the BCS-BEC crossover must also be modified, such that the low-density asymptotic state becomes a gaseous mixture of neutrons and deuterons [24].

Two relevant energy scales for the problem domain under study are provided by the shift $\delta\mu = (\mu_n - \mu_p)/2$ in the chemical potentials μ_n and μ_p of neutrons and protons from their common value μ_0 and the pairing gap Δ_0 in the SD channel at $\delta\mu = 0$. With increasing isospin symmetry, i.e., as $\delta\mu$ increases from zero to values of order Δ_0 , a sequence of unconventional phases may emerge. One of these is a neutron-proton condensate whose Cooper pairs have non-zero center-of-mass (CM) momentum [25, 26]. This phase is the analogue of the Larkin-Ovchinnikov-Fulde-Ferrell (LOFF) phase in electronic superconductors [27, 28]. Another possibility is phase separation into superconducting and normal components, proposed in the context of cold atomic gases [29]. At large isospin asymmetry, where SD pairing is strongly suppressed, a BCS-BEC crossover may also occur in the isotriplet 1S_0 pairing channel. The ideas of unconventional SD pairing and the BCS-BEC crossover in a model of isospin-asymmetric nuclear matter were combined recently in Ref. [21]. A phase diagram for superfluid nuclear matter over wide ranges of density, temperature, and isospin asymmetry was constructed. A self-consistent set of equations, which includes the gap equation and the expressions for the densities of constituents (neutrons and protons) was solved allowing for the ordinary BCS state, its low-

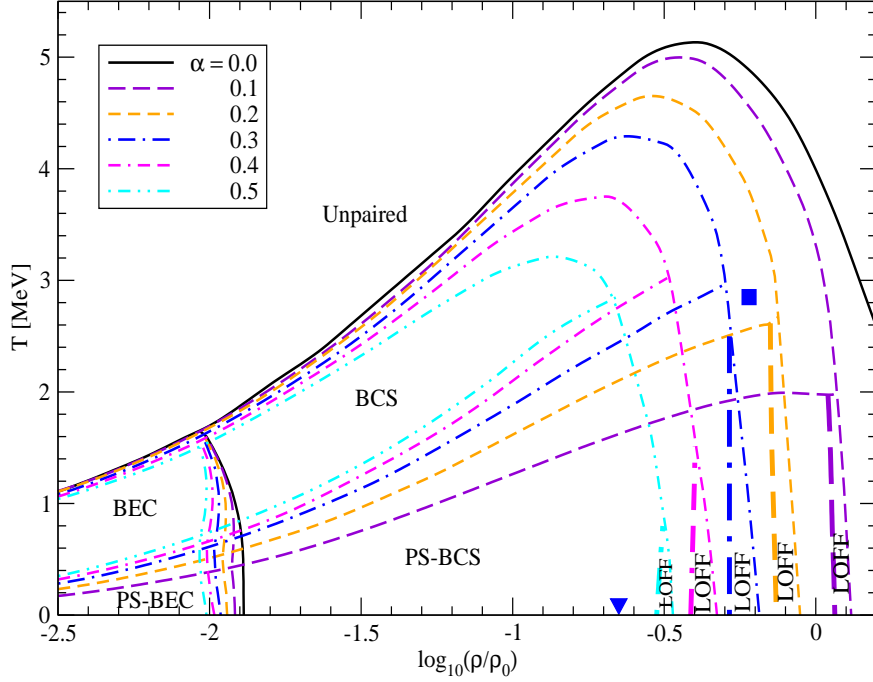


Figure 1. Phase diagram of dilute nuclear matter in the temperature-density plane for several isospin asymmetries α (from Ref. [21]). Included are four phases: unpaired phase, BCS (BEC) phase, LOFF phase, and PS-BCS (PS-BEC) phase. For each asymmetry there are two tri-critical points, one of which is always a Lifshitz point [30]. For special values of asymmetry these two points degenerate into single tetra-critical point at $\log(\rho/\rho_0) = -0.22$ and $T = 2.85$ MeV (shown by a square dot) for $\alpha_4 = 0.255$. The LOFF phase disappears at the point $\log(\rho/\rho_0) = -0.65$ and $T = 0$ (shown by a triangle) for $\alpha = 0.62$.

density asymptotic counterpart BEC state and two phases that owe their existence to the isospin asymmetry: the phase with a moving condensate (LOFF phase) and the phase where the normal fluid and superfluid break down into separate domains.

The phase diagram found in Ref. [21] is shown in Fig 1. The results for the BCS phase and BCS-BEC crossover are consistent with the earlier studies: one observes a smooth crossover to an asymptotic state corresponding to a mixture of a deuteron Bose condensate and a gas of excess neutrons. The transition from BCS to BEC is established according the following criteria: (i) The average chemical potential $\bar{\mu}$ changes its sign from positive to negative values, (ii) the coherence length of a Cooper pair becomes comparable to the interparticle distance, i.e., $\xi \sim d \sim \rho^{-1/3}$ as conditions change from $\xi \gg d$ to $\xi \ll d$.

The nuclear LOFF phase arises as a result of the energetic advantage of translational symmetry breaking by the condensate, in which pairs acquire a non-zero CM momentum \vec{Q} . At constant asymmetry, a temperature increase shifts the gap maximum and the free-energy minimum of the LOFF phase toward small Q , and at sufficiently high temperature and small asymmetry the BCS state is favored over the LOFF phase. This behavior is well understood in terms of the phase-space overlap of the Fermi surfaces of neutrons and protons, which (at finite asymmetry) increases with temperature and the momentum Q of the Cooper pairs.

Thus, as the temperature increases, we expect a restoration of the BCS phase and of the translational symmetry in the superfluid. Obviously, the same restoration occurs when the isospin asymmetry is small enough. The superfluid phase with phase separation (PS) has the

symmetrical BCS phase as one of its components. The temperature dependence of this phase is well established within BCS theory. The second component, which accommodates the neutron excess, is a normal Fermi liquid whose low-temperature thermodynamics is controlled by the excitations in the narrow strip of width $T/\epsilon_{F,n/p}$ around the Fermi surfaces of neutrons and protons.

The transition to the BEC regime of strongly-coupled neutron-proton pairs, which are asymptotically identical with deuterons, occurs at low densities. As already well established, in the case of neutron-proton pairing the criteria for the BCS-BEC transition are fulfilled, i.e., $\bar{\mu}$ changes sign and the mean distance between the pairs becomes larger than the coherence length of the superfluid.

We now turn to the question of how the BCS-BEC crossover is affected by the existence of nuclear LOFF and PS phases at non-zero isospin asymmetries, and conversely how these phases evolve in the strongly-coupled regime if the density of the system is decreased. Four different phases of matter are present in the phase diagram (Fig. 1). (i) The unpaired phase is always the ground state of matter at sufficiently high temperatures $T > T_{c0}$, where $T_{c0}(\rho)$ is the critical temperature of the superfluid phase transition at $\alpha = 0$. (ii) The LOFF phase is the ground state in a narrow temperature-density strip at low temperatures and high densities. (iii) The PS phase appears at low temperatures and low densities, while the isospin-asymmetric BCS phase is the ground state for all densities at intermediate temperatures. In the extreme low-density and strong-coupling regime the BCS superfluid phases have two counterparts: the BCS phase evolves into the BEC phase of deuterons, whereas the PS-BCS phase evolves into the PS-BEC phase, in which the superfluid fraction of matter is a BEC of deuterons. The superfluid-unpaired phase transitions and the phase transitions between the superfluid phases are of second order (thin solid lines in Fig. 1), with the exception of the PS-BCS to LOFF transition, which is of first order (thick solid lines in Fig. 1). The BCS-BEC transition and the PS-BCS to PS-BEC transition are smooth crossovers. At non-zero isospin asymmetry the phase diagram features two tri-critical points where the simpler pairwise phase coexistence terminates and three different phases coexist.

The extreme low-density region of the phase diagram features two crossovers. At intermediate temperatures one recovers the well-known BCS-BEC crossover, where the neutron-proton BCS condensate transforms smoothly into a BEC gas of deuterons with some excess of neutrons. The new ingredient of our phase diagram is the second crossover at low temperatures, where the heterogeneous superfluid phase is replaced by a heterogeneous mixture of a phase containing a deuteron condensate and a phase containing neutron-rich unpaired nuclear matter.

What would be the effect of adding somewhat heavier clusters, notably ^3He , ^3H and ^4He , to the phases discussed above? At low densities statistical equilibrium suggests that the heavier clusters may “eat up” some of the phase-space, therefore their main effect would be to reduce the phase space available to pair-correlated particles, which would eventually lead to some shifts in the phase diagram without qualitative modifications of its structure or topology. The physics might not be as trivial at high densities, for we know that the asymptotically dense phase of matter does not contain clusters. The transition to the homogenous phase can be attributed to the Pauli-blocking of the phase space that can be used to form three and four body bound states. In other words, the rise of the Fermi surface suppresses the formation of three and four-body bound states. It remains an open problem to tackle the interplay of the clusters and pair-correlations on the ultimate composition of matter at subnuclear densities. As stated in the introduction, the answer is not merely of academic interest. The neutrinos (and more generally leptons) transport energy in the supernova processes while moving through this environment; its composition, many-body effects, etc remain the key unknowns in the accurate description of the neutrino transport.

3. Neutrino pair-breaking processes

Next we transport ourselves in time several weeks past the supernova explosion and the formation of a neutron star. During this period of time the star has cooled down to the temperature of superfluid phase transition of baryons, which is roughly in the range $T_c \simeq 0.5$ MeV. At this stage the crust of the star cools predominantly via the electron (e^-) bremsstrahlung process: $e^- + A \rightarrow e^- + A + \nu + \bar{\nu}$, where A stands for a nucleus and ν and $\bar{\nu}$ for the neutrino and antineutrino [31]. The surface of the star cools by emitting thermal soft X -rays (photons). Both processes are not important sinks of energy up until the star's age $t \geq 10^4 - 10^5$ yr [2, 3, 4]. The exact value of the transition to the crust plus surface cooling depends on a number of factors, which we will not discuss. While it was widely accepted from the beginning that the superfluidity will suppress the neutrino radiation processes on baryons in the core, once it becomes superfluid, the less trivial pair-breaking processes were initially neglected, although the rates of these processes, computed at one-loop, were available since 1976 [32, 33, 34]. The Standard Model weak neutral interactions proceed via vector and axial vector interactions. Thus, the neutrino emissivity (phase space integrated rate of neutrino emission) is mainly determined by the response functions of matter to weak vector and axial-vector currents. These can be computed within the finite temperature Green's functions technique in general. Specific applications within the real-time Schwinger-Keldysh formalism can be found in Refs. [33, 35].

Before discussing these response functions, we pause to recall the many-body developments of the formal theory. The polarization tensor in superfluid system involves a re-summation of infinite series of particle-hole diagrams. The methods for doing so were developed in the case of symmetrical nuclear matter, with the purpose of applications to finite nuclei in Ref. [36]. This work is one of the first applications of the Landau Fermi liquid theory (which was designed to describe the properties of non-superfluid liquid He^3) to superfluid systems. The diagrammatic methods of re-summation of particle-hole ladders were developed even earlier [37, 38]. While the theory of superfluid Fermi liquids was generalized further by Leggett [39] to finite temperatures (in the context of condensed matter theories), Leggett did so by computing directly the response functions and not the vertex functions. In nuclear physics, the theory took a turn that focused the entire subsequent work on finite systems (i.e. nuclei), which were treated in the equation of motion formalism for second-quantized operators.

The importance of the vertex corrections for the case of the vector current interactions was first pointed out in Ref. [40], where the authors implemented a polarization tensor adopted from condensed matter work (derived not quite in the regime needed for neutron stars). A number of subsequent works derived directly the finite temperature vertex functions and polarization tensors for neutron and proton components of the core of the star [41, 42, 43, 44]. The current consensus is that the vector current emissivity is suppressed compared to the one-loop results by a factor $(v_f/c)^4$, where $v_f/c \sim 0.1$ is the Fermi velocity of the baryons in units of the speed of light. At the same time it was established that the axial vector interactions are adequately represented by the bare vertices and one-loop polarization tensors. Explicitly, the axial vector emissivity given by [34, 42]

$$\epsilon_\nu = \frac{24G_F^2 g_A^2}{105\pi^3} \nu(0) v_f^2 T^7 I_\nu, \quad I_\nu = z^7 \int_1^\infty dy \frac{y^5}{\sqrt{y^2 - 1}} f_F(zy)^2, \quad (1)$$

where $G_F = 1.166 \times 10^{-5}$ GeV $^{-2}$ is the Fermi weak coupling constant, $g_A = 1.25$ is (unquenched) axial vector coupling constant, $\nu(0)$ is the density of states at the Fermi surface, T is the temperature, $z = \Delta(T)/T$ is the ratio of the pairing gap to the temperature, and $f_F(x) = [1 + \exp(x)]^{-1}$ is the Fermi distribution function. The emissivity of the axial vector current is of the order $(v_f/c)^2$ and, therefore, is the dominant channel of the energy loss.

4. Axion emission from compact stars

The strong sector of the Standard Model may feature a CP violating interaction, which arises due to a topological interaction term in the QCD Lagrangian [45]

$$\mathcal{L}_\theta = \frac{g^2\theta}{32\pi^2} F_{\mu\nu}^a \tilde{F}^{\mu\nu a}, \quad (2)$$

where $F_{\mu\nu}^a = \partial_\mu A_\nu - \partial_\nu A_\mu + gf^{abc}A_{\mu b}A_{\nu c}$ is the gluon field strength tensor, $\tilde{F}_{\mu\nu}^a = \epsilon_{\mu\nu\lambda\rho}F^{\lambda\rho a}/2$, f^{abc} are the structure constants of $SU(3)$ group, θ is the parameter which parametrizes the non-perturbative vacuum states of QCD $|\theta\rangle = \sum_n \exp(-in\theta)|n\rangle$, where n is the winding number characterizing each distinct state of QCD, which is not connected to another by any gauge transformation. The QCD action changes by 2π under the shift $\theta \rightarrow \theta + 2\pi$, i.e., θ is a periodic function with a period of 2π . In presence of quarks the physical parameter is not θ , but $\bar{\theta} = \theta + \arg \det m_q$, where m_q is the matrix of quark masses. Experimentally, the upper bound on the value of this parameter is $\bar{\theta} \leq 10^{-10}$, which is based on the measurements of the electric dipole moment of neutron $d_n < 6.3 \cdot 10^{-26} e \text{ cm}$ [46]. The smallness of $\bar{\theta}$ is the strong CP problem: the Standard Model does not provide any explanation on why this number should not be of order unity.

An elegant solution to the strong CP problem is provided by the Peccei-Quinn mechanism [47, 48, 49]. This solution amounts to introducing a global $U(1)_{PQ}$ symmetry, which adds an additional anomaly term to the QCD action proportional to the axion field a . This term acts as a potential for the axion field and gives rise to an expectation value of the axion field $\langle a \rangle \sim -\bar{\theta}$. The physical axion field is then $a - \langle a \rangle$, so that the undesirable θ term in the action is replaced by the physical axion field. The axion is the Nambu-Goldstone boson of the Peccei-Quinn $U(1)_{PQ}$ symmetry breaking [48, 49], and its effective Lagrangian has the form

$$\mathcal{L}_a = -\frac{1}{2}\partial_\mu a \partial^\mu a + \mathcal{L}_{int}(\partial_\mu a, \psi), \quad (3)$$

where the second term describes the coupling of the axion to fermion fields (ψ) of the Standard Model.

There are ongoing experimental searches for the axion. Cosmology and astrophysics provide strong complementary constraints. Because axions can be effectively produced in the interiors of stars they act as an additional sink of energy. The requirements that the energy loss from a star is consistent with the astrophysical observations place lower bounds on the coupling of axions to the standard model particles, and hence on the Peccei-Quinn symmetry breaking scale [5, 50]. The latter limit translates into an upper limit on the axion mass. Such arguments have been applied to the physics of supernova explosions [51, 52, 53, 54, 55] and white dwarfs [56]. In the case of supernova explosions the dominant energy loss process is the emission of an axion in the nucleon (n) bremsstrahlung $n + n \rightarrow n + n + a$. The same process was considered earlier by Iwamoto as a cooling mechanism for mature neutron stars, i.e., neutron stars with core temperature in the range $10^8 - 10^9$ K [57]. The implications of the axion emission by the modified Urca and nucleon bremsstrahlung, as calculated in Ref. [57], were briefly studied in Ref. [58]. However, as discussed in the previous section the pair-breaking processes are the dominant energy loss channels for temperatures below the critical temperature. Therefore, to obtain bounds on the properties of axions from the cooling of compact stars, the rate of the axion cooling is required. This was computed recently in Ref. [6] in the case of S -wave superfluid.

We now briefly outline this calculation. The energy radiated per unit time in axions (axion emissivity) is given by the phase-space integral over the probability of the emission process

$$\epsilon_a = -f_a^{-2} \int \frac{d^3q}{(2\pi)^3 2\omega} \omega g_B(\omega) \kappa_a, \quad \kappa_a = q_\mu q_\nu \Im \Pi_a^{\mu\nu}(q), \quad (4)$$



Figure 2. The two diagrams contributing to the polarization tensor of baryonic matter, which defines the axion emissivity. The “normal” baryon propagators for particles (holes) are shown by single-arrowed lines directed from left to right (right to left). The double arrowed lines correspond to the “anomalous” propagators with two incoming or outgoing arrows. The horizontal dashed lines represent the axion a .

where q and ω are the axion momentum and energy, $\Pi_a^{\mu\nu}(q)$ is the polarization tensor of the superfluid baryonic matter, $g_B(\omega)$ is the Bose distribution function, f_a is the axion-baryon coupling strength. The polarization tensor of a superfluid obtains contributions from four distinct diagrams that can be formed from the normal and anomalous propagators with four distinct effective vertices [41, 42, 43, 44]. However, for the axial vector perturbations the vertices are not renormalized in the medium and, therefore, one proceeds with the bare vertices, in which case the number of the distinct contributions to the polarization tensor reduces to a sum of two admissible bare loops (see Fig. 2). The axion emissivity obtained from Eq. (4) is given by

$$\epsilon_a = \frac{8}{3\pi} f_a^{-2} \nu(0) v_f^2 T^5 I_a, \quad I_a = z^5 \int_1^\infty dy \frac{y^3}{\sqrt{y^2 - 1}} f_F(zy)^2. \quad (5)$$

The T^5 scaling of the emissivity is understood as follows. The integration over the phase space of neutrons carries a power of T , since for degenerate neutrons the phase-space integrals are confined to a narrow strip around the Fermi surface of thickness T . The axion is emitted thermally and being relativistic contributes a factor T^3 to the emissivity. The one power of T from the energy of the axion and the inverse one power of T from the energy conserving delta function cancel. The transition matrix element is proportional to the combinations of u_p and v_p amplitudes, which are dimensionless, but contain *implicit* temperature dependence due to the temperature dependence of the gap function. This dependence is not manifest in Eq. (5), i.e., was absorbed in the definition of the integral I_a . Thus, the explicit temperature dependence of the axion emission rate Eq. (5) is T^5 . As discussed in Sec. 3 at temperatures of order the critical temperature $T_c \simeq 10^9$ K the superfluid cools primarily by emission of neutrinos via the pair-breaking processes driven by the axial-vector currents (we continue to assume that potential fast cooling via direct Urca processes is prohibited). In a first approximation one may require that the axion luminosity does not exceed the neutrino luminosity, i.e.,

$$\frac{\epsilon_a}{\epsilon_\nu} = \frac{10\pi^2}{f_a^2 G_F^2 g_A^2 \zeta_A} \frac{I_a}{I_\nu} \simeq \frac{59.2}{f_a^2 G_F^2 \Delta(T)^2} r(z) < 1, \quad (6)$$

where $r(z) \equiv z^2(I_a/I_\nu)$. Not far from the critical temperature $\Delta(T) \simeq 3.06 T_c \sqrt{1 - T/T_c}$, which translates into $z = 3.06 t^{-1} \sqrt{1 - t}$, where $t = T/T_c$. Numerical evaluations of the integrals provides the following values $r(0.5) = 0.07$, $r(1) = 0.26$, $r(2) = 0.6$ and asymptotically $r(z) \rightarrow 1$ for $z \gg 1$. Noting that $r(z) \leq 1$, one finally obtains

$$f_a > 5.92 \times 10^9 \text{ GeV} \left[\frac{0.1 \text{ MeV}}{\Delta(T)} \right], \quad (7)$$

which translates into an upper bound on the axion mass

$$m_a = 0.62 \times 10^{-3} \text{ eV} \left(\frac{10^{10} \text{ GeV}}{f_a} \right) \leq 1.05 \times 10^{-3} \text{ eV} \left[\frac{\Delta(T)}{0.1 \text{ MeV}} \right]. \quad (8)$$

The bound Eq. (7) can be written in terms of the critical temperature by noting that $\Delta(T) \simeq T_c$ in the temperature range $0.5 \leq t < 1$ of interest.

The neutrino cooling era of compact stars, which spans the time-period $t \leq 10^4 - 10^5$ yr after their birth in supernova explosions is a sensitive probe of the particle physics of their interiors. If one assumes that there are no rapid channels of cooling in neutron stars, i.e., deconfined quarks, above Urca threshold fractions of protons or hyperons (all of which lead to a rapid Urca cooling), then neutron stars cool primarily by neutrino emission in Cooper pair-breaking processes in baryonic superfluids. If, however, axions exist in Nature, the neutron stars must cool via axion emission in Cooper pair-breaking processes, whose axion emission rate scales as T^5 . This scaling differs from the T^7 scaling of the counterpart neutrino processes. The difference arises from the different phase spaces required for the pairs of neutrinos and the axion and is independent of the baryonic polarization tensor. Note also that the rate of axion emission from a P -wave superfluid will differ from the S -wave rate, derived above, by a factor $O(1)$ and, therefore, will not change quantitatively the obtained bounds on the axion parameters.

Similar bounds to those quoted above were obtained previously by Iwamoto [57] ($f/10^{10}\text{GeV} > 0.3$) from a comparison of the rates of axion bremsstrahlung and modified Urca neutrino emission by mature neutron stars, and by Umeda et al [58] ($f/10^{10}\text{GeV} > 0.1 - 0.2$) from fits of cooling simulations to the PSR 0656+14 data¹. The lower bound on f_a derived in Ref. [6] is somewhat larger than the one that follows from the requirement that the axions do not “drain” too much energy from supernova process so that it fails [5, 51, 52, 53, 54, 55]. Note the dependence of the bound Eq. (7) on the pairing gap. This is an ingredient, which originates from superfluidity of cold neutron stars, that does not appear in other bounds on axion parameters. Because, the magnitude and density dependence of the gap are not well-known (for a review see Ref. [59]) there remains enough room for speculations on the impact of the axions on the cooling of compact stars.

5. Perspectives

Although the equation of state and composition of the dilute and warm nuclear matter have been studied for several decades, there remain a number of unsettled issues related to the complex many-body character of this system. The interplay of the pair-correlations and clustering is one of the challenging problems. As we have seen the phase structure of the matter at finite isospin becomes more complicated due to the emergence of novel superconducting phases. It is very likely that along with the deuteron condensate and its BCS counterpart a BEC of alpha particles will coexist with these phases. Thus, one needs to work out the physics of a mixtures of several superfluids phases in this context. How these new features will affect the neutrino transport in supernovas remains an open question.

Neutron star cooling offers a new playground for the studies of beyond Standard Model physics. The example presented above shows that with microscopic rates of the axion emission from superfluid phases at hand one can place useful bounds on the axion properties. Axions offer a unique channel of rapid energy loss by medium to low mass compact stars, where the central densities are below the direct Urca thresholds and/or below the density of deconfinement into quark matter. Thus, the long-standing paradigm that the light and medium mass neutron stars cannot cool rapidly will not hold should the axions exist in Nature and couple to matter strong enough to cool a compact star faster than the Standard Model neutrinos.

¹ Note that Eq. (9) of this work is incorrect, which could be the reason why the limits on the axion’s mass reported in their Figs. 2-4 are by an order of magnitude larger compared to those quoted in Ref. [6].

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